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Direct and Resolved Pomeron in Rapidity Gap Cross Sections*+

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Abstract

We investigate the effect of a direct pomeron coupling to quarks on inclusive jet production in DIS and photoproduction. The direct pomeron coupling generates a point-like contribution to the diffractive part of the structure function F_2, which is analysed on the basis of the latest H1 and ZEUS data. Our model assumptions for the pomeron structure are consistent with the measured data.

1. Introduction

In diffractive production of hadronic final states in ep scattering, the proton stays intact or becomes a low mass state. Between the direction of the proton remnant, which goes down the beam pipe, and the produced hadronic system there is no colour flow, which allows of the possibility to observe large gaps in rapidity between these directions. The experiments H1 [1] and ZEUS [2] at HERA have measured the portion of diffractive events to be $\approx 10\%$ of all events – not only in photoproduction ($Q^2 < 0.01\,\mathrm{GeV}^2$) but also in deep inelastic scattering (DIS) ($Q^2 > 10\,\mathrm{GeV}^2$).

This paper is organized as follows. First, we describe our ansatz to analyse diffractive ep scattering. In section 3, we consider the pomeron structure function and the diffractive part of F_2. We find consistency with the latest HERA data. Finally, we analyse jet cross sections in diffractive inclusive photoproduction and in DIS.

2. Model for diffractive jet production

There exist various phenomenological models [3, 4, 5, 6] to describe the above mentioned diffractive nonperturbative QCD phenomena quantitatively. We follow a widely spread assumption, where the proton splits off a colourless object called pomeron ($I\!P$), which has the quantum numbers of the vacuum. Then, if factorization holds, the proton vertex can be parametrized by a $I\!P$–flux factor that depends on $t = (p-p')^2$, the momentum transfer to the pomeron, and $x_{I\!P}$, the fraction of proton energy, that it carries away.

In fact, this has been done in the past by various authors, who fixed their parameters with the help of $p\overline{p}$ scattering data. Inspired by Regge phenomenology, Berger et al. [7] found for the pomeron flux

$$f_{I\!P/p}(t, x_{I\!P}) = \frac{\beta^2_{I\!Pp}(t)}{16\pi} x_{I\!P}^{1-2\alpha(t)} \qquad (1)$$

with the Regge trajectory $\alpha = \alpha_0 + \alpha' t$ and the residue

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function $\beta^2_{I\!\!Pp}(t) = \beta^2_{I\!\!Pp}(0)e^{b_0 t}$, where $\alpha_0 = 1+\epsilon$, $\epsilon = 0.085$, $\alpha' = 0.25\,\mathrm{GeV}^{-2}$, $\beta^2_{I\!\!Pp}(0) = 58.74\,\mathrm{GeV}^{-2}$ and $b_0 = 4.0\,\mathrm{GeV}^{-2}$.

This definition of the factorization in a pomeron flux factor differs by a factor of $\frac{\pi}{2}$ from the definition of Donnachie and Landshoff [8]. In addition, they included the Dirac elastic form factor of the proton, which is given by

$$F_1(t) = \frac{(4m_p^2 - 2.8t)}{(4m_p^2 - t)}\frac{1}{(1-t/0.7)^2} \tag{2}$$

and yielded

$$f_{I\!\!P/p}(t, x_{I\!\!P}) \;=\; \frac{9\delta^2}{4\pi^2}[F_1(t)]^2 x_{I\!\!P}^{1-2\alpha(t)} \tag{3}$$

with $\delta = 3.24\,\mathrm{GeV}^{-2}$. Equation 3 may be regarded as the most natural way to define a flux factor [4] and will be used in the following. The same definition of the flux factor was used by Ingelman and Schlein [3]. They made a different parametrization with two exponentials of the following form

$$f_{I\!\!P/p}(t, x_{I\!\!P}) \;=\; \frac{1}{2}\frac{1}{\kappa x_{I\!\!P}}(A\exp^{\alpha t} + B\exp^{\beta' t}) \tag{4}$$

with the parameters $\kappa = 2.3\,\mathrm{GeV}^2$, $A = 6.38$, $\alpha = 8\,\mathrm{GeV}^{-2}$, $B = 0.424$ and $\beta' = 3\,\mathrm{GeV}^{-2}$. Here, the factor of $1/2$ comes in because of normalization to *one* proton vertex.

For our purposes, the momentum transfer $t = (p-p')^2$ to the pomeron has to be integrated out, since the proton remnant is (still) not tagged. We get

$$G_{I\!\!P/p}(x_{I\!\!P}) \;=\; \int_{-\infty}^{t_2} dt\, f_{I\!\!P/p}(t, x_{I\!\!P}) \tag{5}$$

$$\text{with} \qquad t_2 = -m_p^2 x_{I\!\!P}^2/(1-x_{I\!\!P}) \quad .$$

We emphasize that recently published H1 data [1] and ZEUS data [2] confirm this x^{-n} factorization. The actual values for the exponent n are

$$n \;=\; 1.19 \pm 0.06 \pm 0.07 \quad \text{H1 [1]} \quad , \tag{6}$$

$$n \;=\; 1.30 \pm 0.08^{+0.08}_{-0.14} \quad \text{ZEUS [2]} \quad , \tag{7}$$

which are comparable to the exponent $2\alpha(0) - 1 = 1.17$ from Regge analysis. However, since the pomeron might not be a real particle, there could be problems with the interpretation of (5) [4].

Surely, the pomeron is in some sense only a generic object, that serves to parametrize a nonperturbative QCD effect. Although it is not considered as a physical particle that could be produced in the s–channel, we employ the concept of structure functions for it.

For the hadron–like part of the unknown parton density functions of the pomeron, $G_{b/I\!\!P}$, we propose the ansatz

$$\begin{aligned}
\beta\, G_{u/I\!\!P}(\beta) &= \beta\, G_{\overline{u}/I\!\!P}(\beta) = \beta\, G_{d/I\!\!P}(\beta) = \beta\, G_{\overline{d}/I\!\!P}(\beta) \\
&= 6\beta(1-\beta)\frac{1}{5}\frac{1}{1+r} \quad , \qquad (8) \\
\beta\, G_{s/I\!\!P}(\beta) &= \beta\, G_{\overline{s}/I\!\!P}(\beta) = \frac{1}{2}\beta\, G_{u/I\!\!P}(\beta) \quad , \\
\beta\, G_{g/I\!\!P}(\beta) &= 6\beta(1-\beta)\frac{r}{1+r}
\end{aligned}$$

and vanishing charm contributions. Here, $\beta = x/x_{I\!\!P}$ with Bjorken x. These functions obey the sum rule

$$\sum_b \int_0^1 d\beta\, \beta\, G_{b/I\!\!P}(\beta) \;=\; 1 \tag{9}$$

and are, in our case, defined for an input scale of $Q_0^2 = 2.25\,\mathrm{GeV}^2$. Theoretical motivation on the basis of nonperturbative methods can be found in [7] and [9]. The analysis of the authors suggests a behaviour of the diffractive parton distributions between $(1-\beta)^0$ and $(1-\beta)^2$ for $\beta \to 1$.

The parameter r describes the unknown ratio of the gluon to quark content of the pomeron. To get a first insight into the structure of our model, we restrict the number of free parameters and choose a value $r = 3$, which represents simple gluon dominance in the pomeron. We carry out the usual DGLAP–evolution to get the right Q^2 dependence of these functions [10].

Several groups considered the possibility of a direct pomeron coupling to quarks [11, 12, 13, 14]. A direct pomeron coupling corresponds to a δ–function term in the pomeron structure function and produces a leading–twist behaviour in the p_T spectrum. Our purpose is to find criteria that allow us to see a direct pomeron coupling in the data [15].

As a consequence, similarly to $\gamma\gamma$ scattering, the $\gamma I\!\!P \to q\overline{q}$ cross section also contributes to the pomeron structure function. Here, we assume a direct vector coupling of the pomeron to the quarks with coupling strength c. This is not really justified with respect to the C parity. However, the Q^2 dependence, that is $\sim \log Q^2$ at low x, is only weakly dependent on the spin structure. This can be seen, for instance, if one replaces the vector coupling by a scalar one. To remove the collinear singularity, we introduce the regulator quark masses m_q and obtain for the point–like (pl) part

$$\begin{aligned}
&\beta\, G^{pl}_{q/I\!\!P}(\beta, Q^2) \\
&= \beta\,\frac{N_c}{8\pi^2}\, c^2 \Bigg\{ v\left[-1 + 8\beta\,(1-\beta) - \frac{4m_q^2}{Q^2}\beta\,(1-\beta)\right] \\
&+ \left[\beta^2 + (1-\beta)^2 + \frac{4m_q^2}{Q^2}\beta\,(1-3\beta) - \frac{8m_q^4}{Q^4}\beta^2\right] \ln\frac{1+v}{1-v} \Bigg\}
\end{aligned}$$

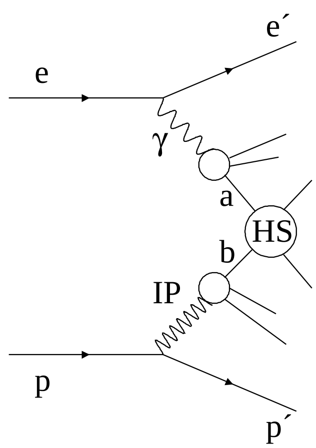

Figure 1. *Generic diagram for the diffractive ep scattering process with a resolved photon γ, a resolved pomeron $I\!\!P$ and the hard subprocess HS.*

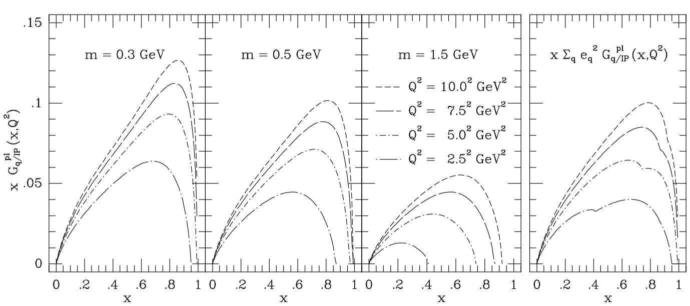

Figure 2. *The pl distribution function for various Q^2 values and our choice of regulator quark masses. The right picture shows the e_q^2 weighted sum that enters into F_2^D. In this plot x is the variable β in (10).*

$$\text{with} \qquad v = \sqrt{1 - \frac{4m_q^2\beta}{Q^2(1-\beta)}} \quad . \qquad (10)$$

The point–like contribution for the three quark masses $m_u = m_d = 0.3\,\mathrm{GeV}$, $m_s = 0.5\,\mathrm{GeV}$ and $m_c = 1.5\,\mathrm{GeV}$ is plotted in figure 2. The right picture shows the e_q^2–weighted sum of the pl contributions from u, d, s and c quarks to F_2^D. The dents in the curves are caused by the charm threshold. At fixed regulator mass, the maximum of the point–like contribution is shifted towards $x = 1$ with increasing scale Q^2.

To get a constraint on the coupling c under the condition $r = 3$, we analyse the recently published data of F_2^D, the diffractive part of the proton structure function.

3. The diffractive contribution to F_2

To leading order in α_s, only the quark distributions of the pomeron enter into the deep–inelastic $I\!\!P$ structure function $F_2^{I\!\!P}(\beta, Q^2)$, which is

$$F_2^{I\!\!P}(\beta,Q^2) = \sum_q e_q^2\,\beta\,\Big[G_{q/I\!\!P}(\beta,Q^2) + G_{\overline{q}/I\!\!P}(\beta,Q^2) + 2\,G^{pl}_{q/I\!\!P}(\beta,Q^2)\Big] \quad . \qquad (11)$$

The comparison with preliminary 1993 H1 data [1] and ZEUS data [2] is shown in figure 3 and figure 4. If factorization holds, which is favoured by the experiments for a large range of β and Q^2 values, the data points are proportional to $F_2^{I\!\!P}(\beta,Q^2)$, i.e. $\tilde{F}_2(\beta,Q^2) = k\,F_2^{I\!\!P}(\beta,Q^2)$ with a constant k that is determined by the $I\!\!P$-flux factor:

$$k = \int_{x_{I\!\!P}^{min}}^{x_{I\!\!P}^{max}} dx_{I\!\!P} \int_{t_1}^{t_2} dt\, f_{I\!\!P/p}(x_{I\!\!P},t) \quad . \qquad (12)$$

The absolute normalization due to k is very sensitive to the integration bounds, $x_{I\!\!P}^{min}$ and $x_{I\!\!P}^{max}$, which are taken from the respective experiments (see figure 3 and figure 4). Unfortunately, no experimental errors on them have been published yet. A small reduction of the integration interval would improve our normalization to the data substantially.

We concentrate therefore on the discussion of the shapes of the data in comparison to our model. We find that our model fits the shape of the data curves well. Especially, the Q^2–evolution is fitted better for a combined ansatz, i.e., quarks in the pomeron *and* pl part (solid curves in figure 3), than for the DGLAP evolved quark distributions (short dashed line in figure 3) or pl part alone. An alternative possibility to fit the data has been represented in [2].

Finally, we fold the pomeron structure function (11) with the pomeron flux factor in eq. (5) to get the diffractive part of the proton deep–inelastic structure function $F_2^D(x,Q^2)$. The relation is

$$F_2^D(x,Q^2) = \int_x^{x_0} dx_{I\!\!P} \int_{t_1}^{t_2} dt\, f_{I\!\!P/p}(x_{I\!\!P},t) F_2^{I\!\!P}(x/x_{I\!\!P},Q^2) \quad . \qquad (13)$$

The upper bound $x_0 = 0.01$ is an experimental choice. In contrast to the analysis of the pomeron structure function, the variable Bjorken x is now fixed (instead of β). In figure 5, we compare with 1993 H1 data [20]. Again, we emphasize the consistency of the choice $c = 1$ ($r = 3$) for the direct pomeron coupling in our model and the data. We use this value in the calculation of the diffractive jet cross sections.

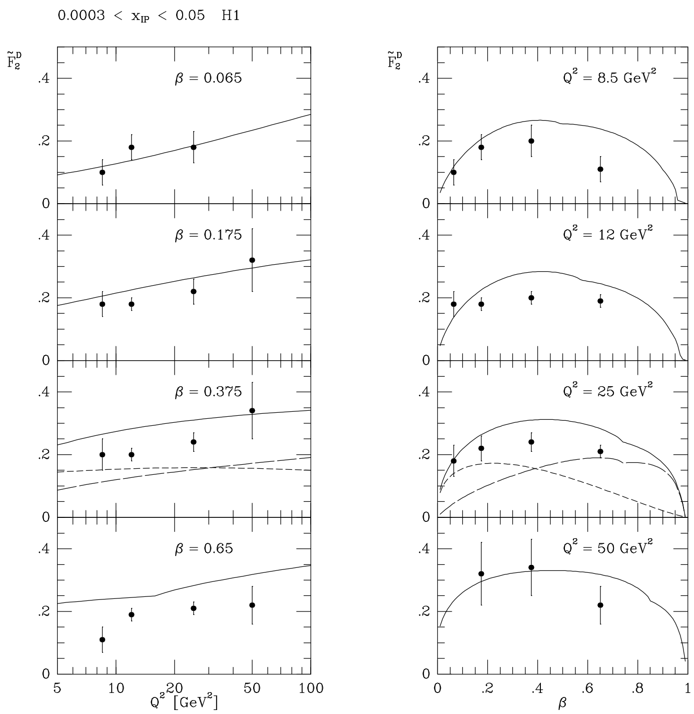

Figure 3. *Comparison of our model predictions with preliminary 1993 H1 data [1]. An important feature of the point–like contribution (long–dashed line) is the filling up of the quark distributions (short–dashed line) at higher β where they become less dominant. This results in flatter curves for the combined distribution (full line) and a better fit to the shape of the data.*

4. Jets in diffractive inclusive photoproduction and DIS

The calculation of the differential jet cross section $\frac{d^2\sigma}{dy dp_T^2}$ of the process depicted in figure 1 is straightforward. Here, p_T and y are the transverse momentum and rapidity of one outgoing jet. In the case of photoproduction, we perform the evaluation in the ep laboratory system where the rapidity is positive for jets travelling in the proton direction. We have then the usual factorization of the photon flux factor at the electron vertex [16]. As is well known in photoproduction, the photon is resolved or couples directly to the final–state quarks. For the photon particle density functions, we take the parametrizations of GRV [17].

A more detailed discussion can be found in [15]. However, here we have included the improved quark distributions of the pomeron due to the performed Q^2–evolution with $Q = p_T$, the transverse momentum of the considered jet. Further, we use in our analysis more actual data of the run in 1993.

The differential inclusive one–jet cross section is

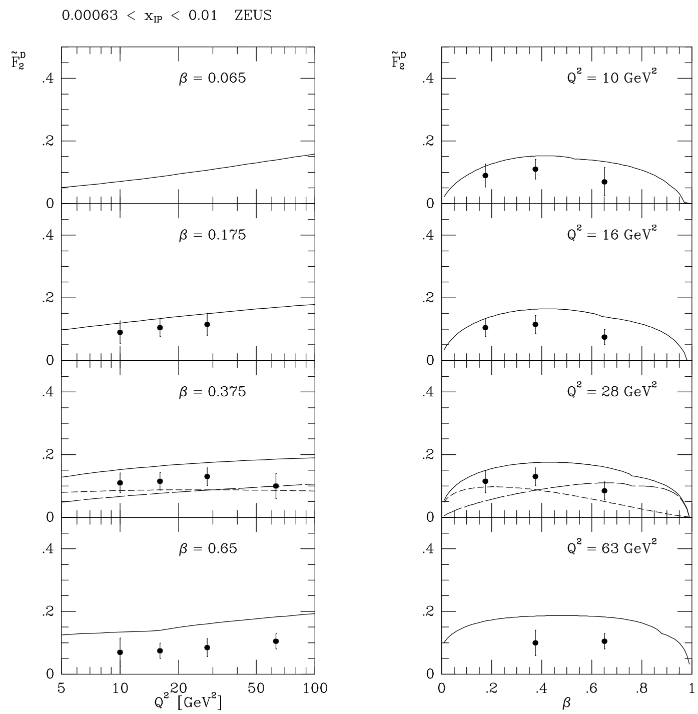

Figure 4. *Comparison of our model predictions with preliminary ZEUS data [2], analogous to figure 3.*

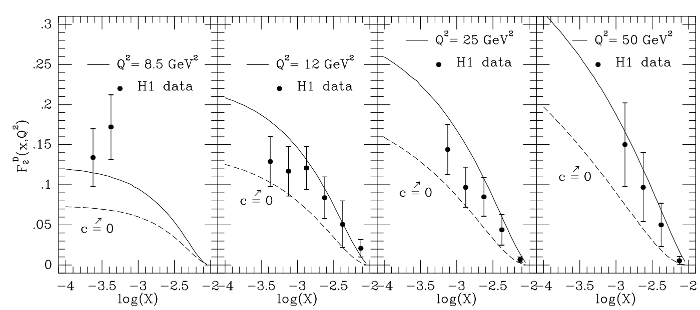

Figure 5. $F_2^D(x, Q^2)$ *compared to preliminary 1993 H1 data [20]. The dashed lines represent only the contributions from quarks in the resolved pomeron, while for the solid curves the pl contribution with $c = 1$ is included.*

obtained by integrating out all kinematic variables over the allowed ranges without regard to the rapidity of the second jet, while for the two–jet cross section, we demand explicitly that the second jet does not enter the cone that is set up by the first jet around the outgoing proton direction.

The results are shown in figure 6. Note the large rapidity gap in the forward direction between the diffractive and nondiffractive parts. In our analysis this is controlled by the $x_{I\!P}^{max} = 0.01$ cut. This

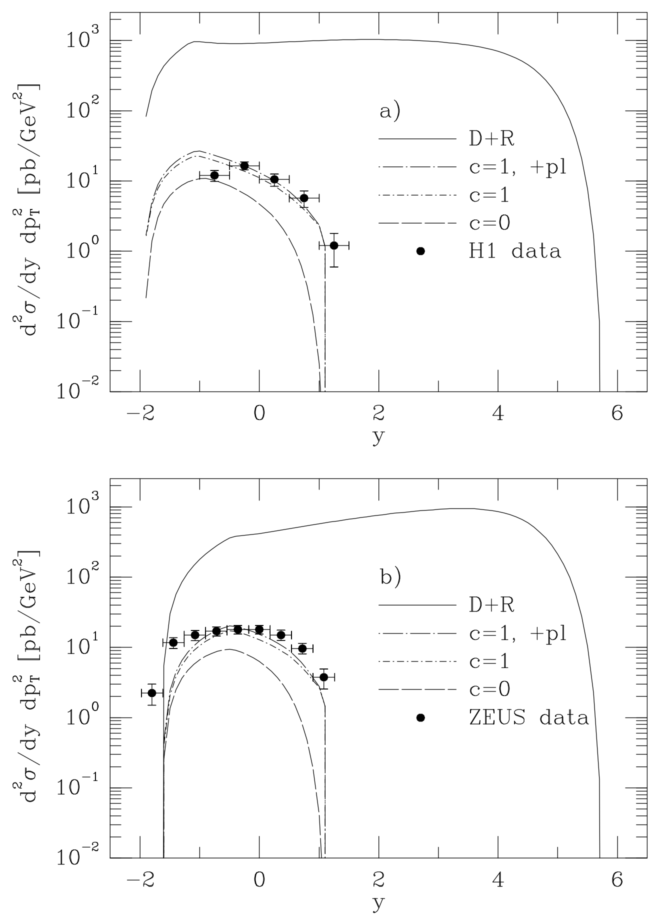

Figure 6. *The rapidity distribution of the a) one– and b) two–jet cross sections for fixed transverse momentum $p_T = 5\,GeV$ in the ep laboratory system. Here, y is defined to be positive for jets travelling in the proton direction. For comparison, the nondiffractive cross section obtained with CTEQ parametrizations of the proton structure functions is also shown (solid line). Since the 1993 event rates (data points) of H1 [21] and ZEUS [22] are not normalized to the luminosity, we can compare only the shapes.*

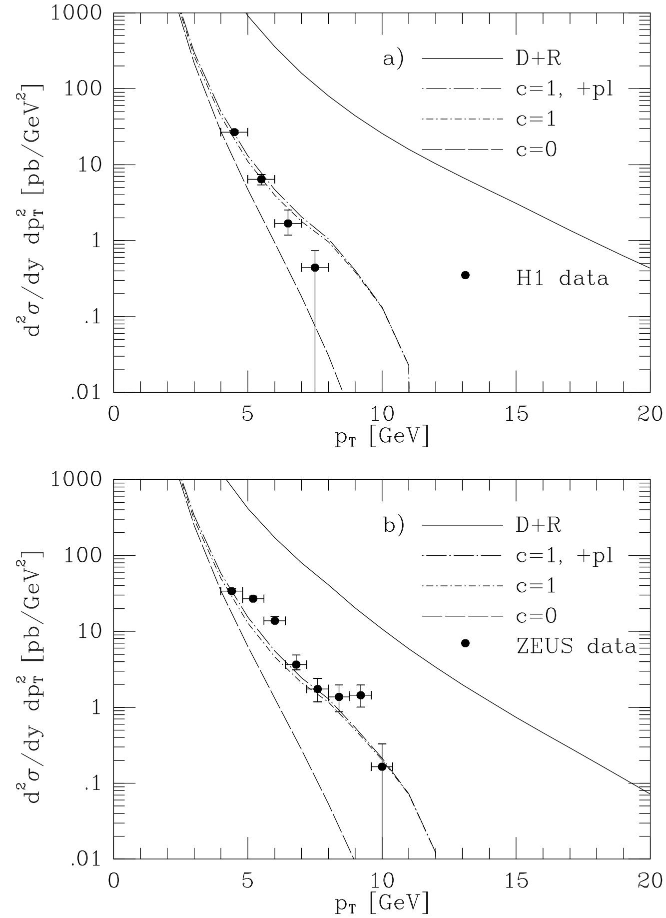

Figure 7. *The p_T spectra of the a) one–jet and b) two–jet cross sections for fixed rapidity $y = 0$. Like in figure 6, the nondiffractive cross section obtained with CTEQ parametrizations of the proton structure functions is also shown (solid line). Since the 1993 event rates (data points) of H1 [23] and ZEUS [22] are not normalized to the luminosity, we can compare only the shapes. For higher p_T values the contribution from a direct pomeron quark coupling dominates the resolved part (c=0) and is favoured by the shape of the data.*

corresponds to a $y^{max} \approx 1.2$ and can be compared to the experimental value $y^{max}_{exp} \approx 1.5$. The lower limit in the y–distribution is due to the cut on x_γ that enters the improved Weizsäcker Williams formula and is numerically determined by the experimental conditions.

The need of a direct pomeron coupling becomes clear, if one inspects the slope of the p_T spectra in figure 7. As expected, with the direct coupling, the p_T spectrum does not fall off so strongly compared to the resolved pomeron contribution ($c = 0$). The point–like component of the structure function, however, does not play a significant role in this discussion, which can be understood by the gluon dominance in our model ($r = 3$ in eqs. (8)).

A reduction of r would increase the quark content in the pomeron due to the sum rule, eq. (9). But this would be accompanied by a reduction of the coupling c to satisfy the bounds coming from the analysis of the diffractive part of F_2 in section 3.

In DIS, the photon is always direct. The jet cross section has been calculated for our model in the $\gamma^* p$ cms. Jets with positive rapidity are travelling now in the photon direction. With p, the four–momentum of the proton, k the four–momentum of the electron, and q, the momentum transfer to the photon, we define the usual kinematic variables $s = (p+k)^2$, $Q^2 = -q^2$, $W^2 = (p+q)^2$, $x = Q^2/(2pq)$, $y_e = pq/pk$, $z = p_T e^y/W$. Finally, we set $\xi = x + p_T^2/(yz(1-z)s)$, which is the fraction of energy delivered from the proton to the subprocess (see figure 1).

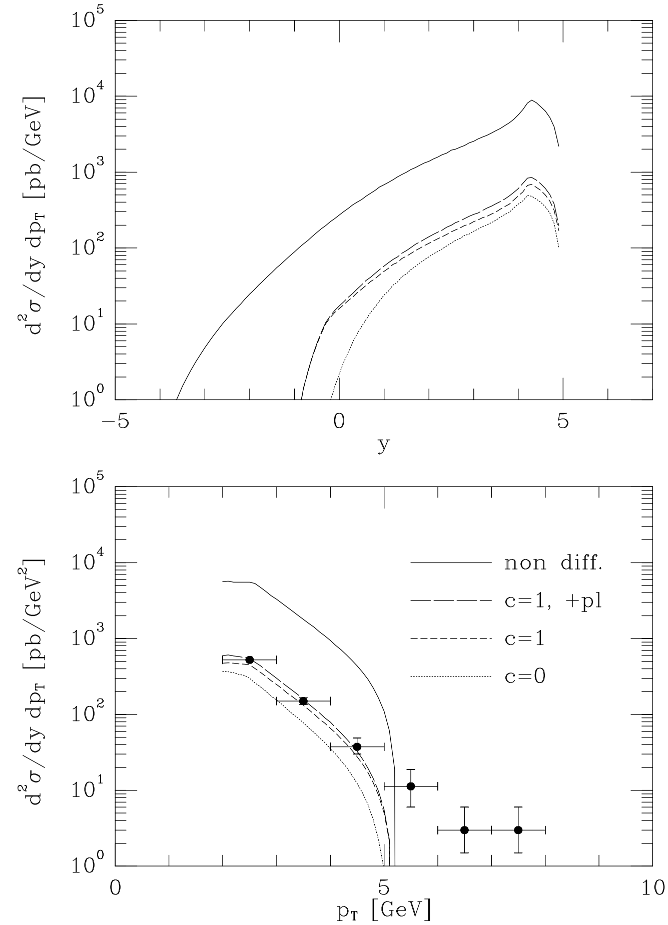

Figure 8. *The y distribution and p_T spectrum of the one jet cross section in DIS, evaluated in the $\gamma^* p$ cms for fixed $p_T = 2\,GeV$ and $y = 4$, respectibly. We impose the experimental conditions $Q^2 > 10\,GeV^2$ and $W^2 > 140^2\,GeV^2$ which correspond to the data points from ZEUS [24].*

The inclusive one–jet cross section is then given by

$$\frac{d^2\sigma}{dy\,dp_T} = \int_{a+b}^{1} dy_e \int_{a/y_e}^{1-b/y_e} dx \frac{d^4\sigma}{dx\,dy_e\,dy\,dp_T} \quad ,(14)$$

where the kinematic bounds follow from the requirements $Q^2 \geq Q^2_{min}$ and $W^2 \geq W^2_{min}$: $a = Q^2_{min}/s$, $b = \max((2p_T \cosh y)^2, W^2_{min})/s$.

In the resolved pomeron case, we have

$$\frac{d^4\sigma}{dx\,dy_e\,dy\,dp_T^2} \tag{15}$$
$$= \sum_{bi} \int_{\xi}^{x_{I\!P}^{max}} \frac{dx_{I\!P}}{x_{I\!P}} G_{I\!P/p}(x_{I\!P}) G_{b/I\!P}(\xi/x_{I\!P}, Q^2)$$
$$\frac{2\alpha_s \alpha^2 Q_i^2}{\xi x y_e^3 (1-z) s^2} \left\{ [1 + (1-y_e)^2] \hat{h}_u + 2(1-y_e)\hat{h}_l \right\} \quad .$$

The contribution of the direct coupling is

$$\frac{d^4\sigma}{dx\,dy_e\,dy\,dp_T^2} \tag{16}$$
$$= \sum_i G_{I\!P/p}(x_{I\!P}) \frac{2c^2\alpha^2 Q_i^2}{4\pi\xi x y_e^3 (1-z) s^2} \tag{17}$$
$$6 \left\{ [1 + (1-y_e)^2] \hat{h}_u + 2(1-y_e)\hat{h}_l \right\} \quad .$$

The functions $\hat{h}_u \equiv \frac{1}{2}(\hat{h}_g + \hat{h}_l)$ and $\hat{h}_l$ depend on the Mandelstam variables $\hat{s} + \hat{t} + \hat{u} = -Q^2$ of the subprocesses. For $\gamma q \to qg$, we have

$$\hat{h}_g = \frac{4}{3}\left(-\frac{\hat{t}}{\hat{s}} - \frac{\hat{s}}{\hat{t}} + \frac{2Q^2\hat{u}}{\hat{s}\hat{t}} \right) \quad , \tag{18}$$
$$\hat{h}_l = \frac{4}{3} \frac{-2Q^2\hat{t}}{(Q^2+\hat{s})^2} \quad , \tag{19}$$

while for $\gamma g \to q\overline{q}$, we find

$$\hat{h}_g = \frac{1}{2}\left(\frac{\hat{t}}{\hat{u}} + \frac{\hat{u}}{\hat{t}} - \frac{2Q^2\hat{s}}{\hat{t}\hat{u}} \right) \quad , \tag{20}$$
$$\hat{h}_l = \frac{1}{2} \frac{4Q^2\hat{s}}{(Q^2+\hat{s})^2} \quad . \tag{21}$$

Like in the photoproduction case, absolute experimental data for the rapidity distribution or p_T spectrum are not yet available to us. In figure 8, we compare the shape of the p_T spectrum with 1993 ZEUS data [24]. The experimental conditions were $Q^2 > 10\,GeV^2$, $W > 140\,GeV$ and $0.04 < y_e < 0.95$. No special case of our model is favoured,since the slopes of direct and resolved contribution are identical, but the shape can be approximately reproduced.

In our analysis, we did not consider particle production or hadronization effects etc. The Monte–Carlo–programs POMPYT by Bruni and Ingelman [18] or RAPGAP by Jung [19] have been developed for a wide class of pomeron models and allow, together with other programs, the study of event characteristics. They are widely used by the HERA collaborations to interpret the large rapidity gap data.

5. Conclusion

In summary, we have studied the effect of a direct pomeron coupling on diffractive jet production at HERA. The concept of pomeron structure functions with DGLAP Q^2–evolution has to be enlarged if the pomeron has a direct coupling to quarks. We have included the additional point–like part in the analysis of diffractive F_2 data and find consistency for the assumption of a direct pomeron coupling to quarks. Some evidence for a direct coupling has been found in the p_T spectrum of photoproduction. However,

our analysis is model dependent and second, except for the discussion of $\tilde{F}_2^D$ and F_2^D, we have compared only the shapes and not the normalizations of the photoproduction and DIS cross sections.

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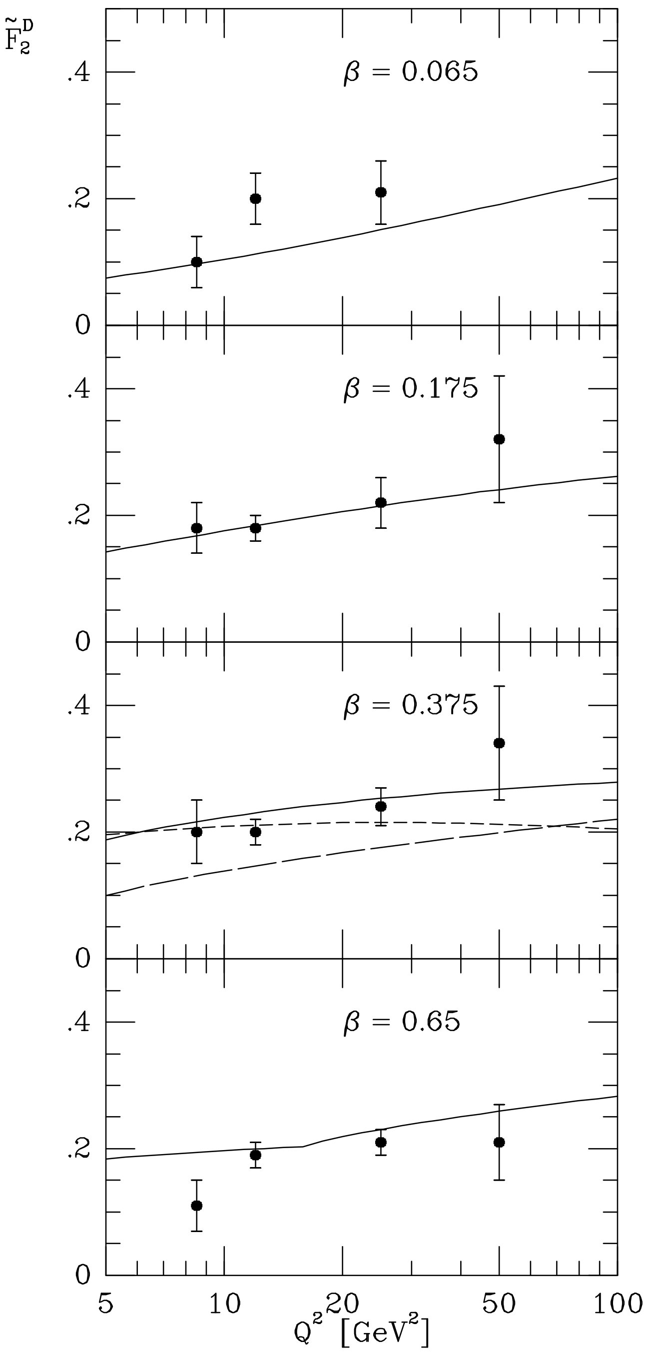

$\tilde{F}_2^D$
$\beta = 0.065$
$\beta = 0.175$
$\beta = 0.375$
$\beta = 0.65$
Q^2 [GeV^2]
$Q^2 = 8.5$ GeV^2
$Q^2 = 12$ GeV^2
$Q^2 = 25$ GeV^2
$Q^2 = 50$ GeV^2
β